

# Post-inflationary brane cosmology

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The brane cosmology has invoked new challenges to the usual Big Bang cosmology. In this paper we present a brief account on thermal history of the post-inflationary brane cosmology. We have realized that it is not obvious that the post-inflationary brane cosmology would always deviate from the standard Big Bang cosmology. However, if it deviates some stringent conditions on the brane tension are to be satisfied. In this regard we study various implications on gravitino production and its abundance. We discuss Affleck-Dine mechanism for baryogenesis and make some comments on moduli and dilaton problems in this context.

## I. INTRODUCTION

Recently, there has been a great deal of interest in conceiving our universe to be a brane-world embedded in a higher dimensional space-time [1]. Such a claim has a motivation from strongly coupled string theories [2,3]. The field theory limit of strongly coupled  $E_8 \times E_8$  heterotic string theory (or M-theory) is believed to be an 11 dimensional supergravity theory [2]. The 11 dimensional world comprises two 10 dimensional hypersurfaces embedded on an orbifold fixed points. The fields are assumed to be confined to these hypersurfaces, known as 9 branes, which can be seen as forming the boundaries of the space-time. It has also been shown that after compactifying the 11 dimensional field theory on a Calabi-Yau three fold it is possible to obtain an effective 5 dimensional theory [4]. The five dimensional space has a structure  $\mathcal{M}_5 = \mathcal{M}_4 \times S^1/Z_2$  and contains two three dimensional branes situated on the orbifold fixed points. The structure allows  $\mathcal{N} = 1$  gauged supergravity in the bulk and on the orbifold fixed planes which could be a realistic testable model for a particle physics phenomenology.

One the other hand there has been some interesting proposals to solve the hierarchy between the two apparent scales the Planck and the electro-weak by introducing extra dimensions and also recognizing the higher dimensional Planck mass to be the fundamental scale. This set-up does not require the world to be supersymmetric. However, this requires some stringent conditions upon gravity residing in the brane and the bulk while the standard model fields are stuck to the observable world [5]. Then there came another twist in constructing models with extra dimensions. It has been shown that gravity can also be localized [6]. The authors have demonstrated that in a background of a special non-factorizable geometry an exponential warp factor appears to the Poincaré invariant 3+1 dimensions. The model consists of two 3 branes situated rigidly along the 5th dimension compactified on a  $S^1/Z_2$  orbifold symmetry. The space-time in the bulk is 5 dimensional anti de-Sitter space and the two branes have opposite brane tensions.

All these new ideas invoke a great concern to the cosmological evolution of the Universe and in this regard

several authors have studied this question with a great emphasis on inflationary cosmology. It has also been noticed that the two hypersurfaces with opposite brane tensions behave differently and they have different cosmology all together [7]. The most startling result was noticed to be departure from the usual four dimensional evolution equation on the brane [10,13]. The presence of branes and the requirement that the fields are localized to the respective branes lead to a non-conventional brane cosmology, which requires lot more study. In this paper we concentrate upon one of the most important aspects of post-inflationary cosmology in our brane.

The idea of cosmological Inflation has many virtues, solving a range of otherwise troubling problems. However, inflation leads to extremely cold Universe with entropy not sufficient to synthesize light elements. Hence, the Universe has to be reheated upto a temperature sufficient enough to have nucleosynthesis, which is close to  $\mathcal{O}(\text{MeV})$ . This is attained via the decay of coherent oscillations of the scalar field whose potential energy dominated the Universe before their decay. The Universe ultimately reaches a thermal equilibrium and also radiation dominated era from where on it follows the standard picture of Big Bang cosmology. The maximum temperature attained during the radiation dominated era is known to be the reheat temperature and to match the bound coming from nucleosynthesis it should be at least more than  $\mathcal{O}(\text{MeV})$ . The reheat temperature plays an important role in the standard Big Bang cosmology and depending on the efficiency of reheating, the reheat temperature could be higher or lower, and in either scenarios there could be interesting cosmological impacts. Interesting results have been obtained for a Universe with a low reheat temperature, see Ref. [8].

In this paper we will be estimating the reheat temperature and then discussing various implications on our brane cosmology. Strictly speaking we will be treating the observable brane as a hypersurface. We will be assuming that at least in our brane supersymmetry is required to solve the hierarchy problem. In this regard we are closer to the string scenario where the effective four dimensional Lagrangian accommodates  $\mathcal{N} = 1$  supergravity with chiral and gauge multiplets. However, to

keep the discussion quite general we will not a priori fix the volume of the six dimensional manifold or the length of the eleventh dimensional segment. This will give us ample choice to the five dimensional Planck mass, which we denote here by  $M_5$ . We begin with a short introduction to chaotic inflation. We will then discuss reheating temperature, gravitino abundance, and towards the end we will briefly discuss the viability of Affleck-Dine mechanism for baryogenesis. We then conclude our paper with some discussions.

## II. CHAOTIC INFLATION ON THE BRANE

It has been noticed in Refs. [9,10] in the context of extra dimensions and the brane-world scenario that the effective four dimensional cosmology is non-trivial and could possibly deviate from a simple Big Bang cosmology. Such a claim has a motivation from string theory which perhaps could lead us completely different physics to the early Universe. It is thus interesting to study the consequences of string motivated early cosmology. However, due to advancement of the observational cosmology, it is no longer believed that the early cosmology is not well constrained. In this regard inflation which is still one of the best paradigms of the early Universe is also constrained by the COBE data which has measured the temperature of the cosmic microwave background radiation and observed a small inhomogeneity which is one part in  $10^5$ . Thus it would be interesting to describe inflation in four dimensions in the context of brane-world scenario. However, the noted deviation for the Friedmann equation at very low temperatures leads to interesting consequences which we will be discussing in this section. Followed by the discussions in Refs. [10,13] we notice that the presence of an extra dimension,  $y$ , compactified on an orbifold  $y = -y$  leads to extra terms in the Friedmann equation in the observable brane. We mention here only the leading terms

$$H^2 = \frac{8\pi}{3M_p^2} \rho \left[ 1 + \frac{\rho}{2\lambda} \right] + \frac{\Lambda_4}{3}, \quad (1)$$

where  $\Lambda_4$  is a four dimensional cosmological constant,  $\rho$  is the energy density of the matter stuck to the brane. In our discussion we will be assuming the four dimensional cosmological constant to be precisely zero from the onset of inflation. The brane tension  $\lambda$  relates the four dimensional Planck mass to the five dimensional Planck mass via

$$M_p = \sqrt{\frac{3}{4\pi}} \left( \frac{M_5^2}{\sqrt{\lambda}} \right) M_5. \quad (2)$$

It is evident that Eq. (1) leads to the usual relation  $H = \sqrt{8\pi\rho/3M_p^2}$  \* when  $\rho < 2\lambda$ . During nucleosynthesis in our brane this is precisely the criteria to be followed because at that time the expansion rate is determined by the energy density linear in  $\rho$  in Eq. (1). This leads to constraining the brane tension  $\lambda > (1\text{MeV})^4$ , and also the five dimensional Planck mass  $10^4\text{GeV} < M_5$ . However, as we shall soon notice that the upper bound on the five dimensional Planck mass will also be fixed by demanding that the inflaton field would not take a value more than the Planck scale in four dimensions during inflation. Here we briefly discuss some of the perspectives of the massive inflaton field with a potential  $V = (m_\phi^2 \phi^2/2)$ . It is important to mention that the presence of an extra dimension does not alter the conservation equation for the matter field stuck to the observable brane.

$$\dot{\rho} + 3H(\rho + p) = 0. \quad (3)$$

This is not always true, especially if there were a scalar field non-minimally coupled to gravity living in the brane and the bulk, then the local conservation equation for a matter field stuck to the brane would not hold. Interesting cosmology will be discussed elsewhere, but from now on we follow Eq. (3). This has an obvious consequence to the scalar field dominating the early Universe during its potential dominated phase. The dominance of  $\rho^2$  term in Eq. (1) leads to enhancing the Hubble friction term in Eq. (3). This naturally assists inflation provided a scalar field is slowly evolving on a potential. It has been demonstrated in various Refs. [11,14,15] that chaotic inflation is possible in this non-conventional scenario. However, to satisfy the COBE data on the observed density perturbations it is pertinent to constrain the mass of the inflaton field which is given by [11]

$$m_\phi \approx 5 \times 10^{-5} M_5 \equiv 4 \times 10^{-5} M_p^{1/3} \lambda^{1/6}, \quad (4)$$

and the inflaton field [11]

$$\phi_{\text{cobe}} \approx 3 \times 10^2 M_5 \equiv 2 \times 10^2 M_p^{1/3} \lambda^{1/6}, \quad (5)$$

where  $\phi_{\text{cobe}}$  is determined by the number of e-foldings required for generating an adequate density perturbations. While deriving the last equation the authors in Ref. [11] have taken  $N_{\text{cobe}} \approx 55$  e-foldings, and, we remind ourselves that the above bounds have been obtained while assuming that the Hubble parameter is dominated by  $\rho^2$  term in Eq. (1). If we do not want to plague the inflaton potential by non-renormalizable quantum corrections, we would require to begin inflation at a scale below the four dimensional Planck mass, which then allows  $M_5 < 10^{17}$

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\*We will frequently imply Eq. (1) to be a consequence of non-conventional brane cosmology compared to the standard cosmology where  $H = \sqrt{8\pi\rho/3M_p^2}$ .

GeV, thus constraining the value of the five dimensional Planck mass to be  $10^{17}\text{GeV} < M_5 < 10^4\text{GeV}$ , and similarly constraining the brane tension  $10^{64}(\text{GeV})^4 < \lambda < 1(\text{MeV})^4$ .

In this regard we notice that if we assume that our brane has  $\mathcal{N} = 1$  supersymmetry and it is broken at a suitable scale in a hidden sector which is mediated to the observable sector only through gravitational interactions such that it generates soft supersymmetry breaking masses to the gravitino, sfermions and the gauginos, then to solve the gauge hierarchy problem the gravitino mass should be around  $\mathcal{O}(\text{TeV})$ . If we also believe that supersymmetry in our world is a consequence of low energy string theory beyond supersymmetric standard model, then we get dilaton and moduli supermultiplets which will acquire a small mass  $\sim \mathcal{O}(m_{3/2})$ . It has been noted that their cosmology is very similar to the gravitinos [45,46]. Especially, the gravitino abundance is strongly connected to the reheat temperature, because they can be created after inflation. We will be studying them in the next section.

### III. REHEAT TEMPERATURE

One of the virtues of inflation is that the inflaton after the end of inflation becomes extremely homogeneous and begins coherent oscillations around the bottom of the potential. For a massive inflaton field the average pressure vanishes during the oscillations and the energy density of the inflaton follows  $\rho_\phi \propto a^{-3}$ , where  $a$  is the scale factor. If the energy density of the decaying inflaton is larger than the brane tension, then the Hubble expansion as a function of the scale factor can be written as

$$H^2(a) \approx \frac{8\pi}{3M_{\text{P}}^2} \frac{\rho_{\phi i}^2}{2\lambda} \left(\frac{a_i}{a}\right)^6, \quad (6)$$

where we have denoted  $\rho_{\phi i}$  and  $a_i$  as the inflaton energy density and the scale factor at the beginning of the coherent oscillations. Depending on the decay rate of the inflaton the reheating process could be efficient or inefficient. The only observed fact is that the reheat temperature should be more than  $\sim \mathcal{O}(\text{MeV})$  to pave a successful nucleosynthesis. Equating Eq. (6) to the decay rate  $\Gamma_\phi$ , and then equating  $\rho_\phi$  to the energy density of the relativistic species  $\rho_r = (\pi^2/30)g_*T_{\text{rh}}^4$ , where  $g_*$  is the relativistic degrees of freedom, we obtain the reheat temperature of the Universe [18]

$$T_{\text{rh}} \approx \left(\frac{\Gamma_\phi M_{\text{P}} \sqrt{\lambda}}{g_*}\right)^{1/4} \approx \left(\frac{\Gamma_\phi M_5^3}{g_*}\right)^{1/4}. \quad (7)$$

For renormalizable couplings to the inflaton, the decay rate of the inflaton is estimated to be  $\Gamma_\phi \sim \alpha_\phi m_\phi$ , where  $\alpha_\phi$  is a dimensionless Yukawa coupling. The estimation of reheat temperature is given by [18]

$$\begin{aligned} T_{\text{rh1}} &\approx 10^{-5/4} M_{\text{P}}^{1/3} \lambda^{1/6} \left(\frac{\alpha_\phi}{g_*}\right)^{1/4}, \\ &\approx 10^{-5/4} M_5 \left(\frac{\alpha_\phi}{g_*}\right)^{1/4}. \end{aligned} \quad (8)$$

The inflaton field could also decay via gravitational interactions such as in the case of supergravity the decay rate is given by  $\Gamma_\phi \sim m_\phi^3/M_{\text{P}}^2$ . Then the estimation of reheat temperature is given by

$$T_{\text{rh2}} \approx 10^{-3} \left(\frac{\lambda}{g_*}\right)^{1/4} \approx 10^{-3} \frac{M_5^{3/2}}{g_*^{1/4} M_{\text{P}}^{1/2}}. \quad (9)$$

Thus in principle one can find a range of reheat temperatures for these two cases we have discussed, but before analysing them it is crucial to ensure whether the energy density of the thermal bath is more than the brane tension, otherwise our naive assumption behind estimating the reheat temperature, namely Eqs. (8) and (9) would be of no use. So, let us analyse the second scenario when the inflaton was decaying via gravitational interactions. It is not difficult to realise that  $\rho_r(T_{\text{rh2}}) \approx 10^{-11}\lambda$ . This tells us that we were wrong behind our naive assumption. The conclusion is very simple and it suggests that if the inflaton is decaying via gravitational interactions, whatsoever be the value of the brane tension, after inflation we would always be in a regime where  $\rho_r < \lambda$ , and, thus we can safely assume the standard Big Bang lore  $H \propto \sqrt{\rho}/M_{\text{P}}$  instead of a non-conventional brane cosmology. This has an interesting implication that if we were to decide a swift transition from non-conventional inflationary cosmology to the standard Big Bang cosmology, then perhaps this could be easily achieved via slow decaying of the inflaton field.

Now let us similarly analyse the first scenario when the inflaton field was decaying via Yukawa couplings. Thus to ensure that  $\rho_{\text{rh}}$  is greater than  $\lambda$ , we get the following inequality

$$\rho_r(T_{\text{rh1}}) \equiv \frac{\alpha_\phi}{3} 10^{-5} M_5^4 > \lambda \equiv \frac{3}{4\pi} \frac{M_5^6}{M_{\text{P}}^2}, \quad (10)$$

which will be true when the five dimensional Planck mass is constrained by

$$M_5 < \frac{\sqrt{4\pi\alpha_\phi}}{10^3} M_{\text{P}}, \quad (11)$$

which can be amply satisfied if we lower the five dimensional Planck mass. It is noticeable that the upper bound on the five dimensional Planck scale is at least an order of magnitude lower than the bound obtaining from beginning inflation below four dimensional Planck scale. This tells us an important message that if Eq. (11) is satisfied, then we may safely assume Eq. (6), provided that the inflaton field is mainly decaying via the Yukawa couplings rather than pure gravitational couplings. This is an important conclusion which we have to bear in mind. From

here onwards we will only concentrate upon inflaton decaying via Yukawa couplings which are not Planck mass suppressed. Only while discussing Affleck-Dine baryogenesis in non-conventional cosmology we will be assuming that the inflaton decaying via gravitational coupling.

Since we know that the transition from non-conventional cosmology to the standard cosmology should take place before nucleosynthesis, we must then estimate the radiation temperature when the transition between non-conventional to the standard Big Bang cosmology takes place. If we naively assume that the transition is instantaneous, then the transition temperature is given by

$$T_{\text{transit}} \approx \frac{M_5^{3/2}}{M_{\text{p}}^{1/2}}, \quad (12)$$

where we have assumed  $g_* \sim 300$ .<sup>†</sup> It is evident from Eq. (12) that higher the five dimensional Planck mass  $M_5$  is, swifter the transition is. For instance if  $M_5 \sim 10^{16}$  GeV, the transition from non-conventional to conventional cosmology takes place when  $T_{\text{transit}} \approx 10^{14}$  GeV. It is also important to notice that this happens very close to the reheat temperature. However, for a low five dimensional Planck mass such as  $M_5 \sim 10^6$  GeV, the transition temperature could be as low as  $T_{\text{transit}} \leq 1$  GeV. Notice that this happens after the reheat temperature  $T_{\text{rh}} \approx 10^4$  GeV.

Now we can estimate the range of reheat temperature but only for a case when the inflaton is decaying via Yukawa couplings.

$$10^{16} \left( \frac{\alpha_\phi}{g_*} \right)^{1/4} \text{ GeV} \geq T_{\text{rh}} \geq 10^3 \left( \frac{\alpha_\phi}{g_*} \right)^{1/4} \text{ GeV}. \quad (13)$$

We notice that depending on the brane tension, or, the Planck mass in five dimensions, we can have completely different scenarios in the energy scale beyond nucleosynthesis. In this paper we will be discussing some of the concerning issues later on. So far we have been assuming that the reheating would be prompt, or, at least the inflaton would decay via the Yukawa couplings, or, the gauge couplings. However, if the reheating was not very prompt then one could imagine that during the process of reheating the Universe could have a different temperature and perhaps more than the reheat temperature of the Universe. In fact the suspicion is correct, and indeed the Universe can achieve a temperature higher than the reheat temperature. This is due to the fact that if reheating was not prompt, then there could be a short spell

where apart from the decaying inflaton, there could be a radiation content and also the decay products of the inflaton. The equation of state of the matter content would neither mimic radiation nor a pressureless fluid, rather completely different. It has been noticed that during this time in the standard cosmology the temperature follows  $T \propto a^{-3/8}$  [12]. Naively, one would expect similar kind of behaviour in the non-conventional case as well. However, it is not very difficult to realize that following the arguments in Ref. [12], the maximum temperature achieved in non-conventional cosmology is  $T_{\text{MAX}} \approx \mathcal{O}(1)T_{\text{rh}}$ . It can also be noticed that the abundance of a stable massive particles are different in this scenario and work in this direction is in progress. In the next section we describe the gravitino abundance in non-conventional cosmology.

#### IV. GRAVITINO PRODUCTION AND THEIR ABUNDANCE

If we imagine that supersymmetry plays an important role in the early Universe then our scope of discussion enhances a lot, and, out of which the gravitino production and their abundance gets most of our attention. In the gravity mediated supersymmetry breaking the gravitino gets a mass around  $\mathcal{O}(\text{TeV})$ . Since their couplings to other particles are Planck mass suppressed, the life time of gravitino at rest is quite long  $\tau_{3/2} \sim M_{\text{p}}^2/m_{3/2}^3 \sim 10^5(m_{3/2}/\text{TeV})^{-3} \text{ sec}$  [19]. We know that successful nucleosynthesis depends on the ratio of the number density of baryons to photons. The gravitino decay products could easily change this ratio. Their decay products such as gauge bosons and its gaugino partners, or, high energy photons could generate a large entropy which would heat up the photons compared to  $\tau$  and  $\mu$  neutrinos. The abundance of neutrinos essentially determines the  $^4\text{He}$  abundance, and this way even if the gravitinos are not the lightest supersymmetric particles they could cause considerable harm. It was first pointed out in Ref. [17] that the gravitino mass should be larger than  $\sim 10$  TeV in order to keep the successes of the Big Bang nucleosynthesis. On contrary if the gravitinos were stable and if their mass exceeded 1 KeV, they could easily overclose the Universe in absence of inflation [16]. Thus studying gravitino abundance is a paramount in this regard. Especially if their energy density is more than the brane tension, then some interesting physics may take place which were not present in the standard cosmology. With this hope we explore the gravitino abundance in this section.

This section will be reviewed from the results obtained in Ref. [18]. The gravitinos can be created in a thermal bath. For helicity  $\pm 3/2$  gravitinos the cross-section is given by  $\sigma \propto (g^2/M_{\text{p}}^2)$ , where  $g$  is the gauge coupling constant and for helicity  $\pm 1/2$  gravitinos the cross-section is given by  $\sigma \propto (g^2 m_g^2/M_{\text{p}}^2 m_{3/2}^2)$  [20]. In either case we notice that the cross-section is suppressed by

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<sup>†</sup>For illustrating purposes we choose  $\alpha_\phi \sim 0.1$  and  $g_* \sim 300$  at  $T = T_{\text{rh}}$ . However, this would also hold as long as the temperature is greater than the masses of the superpartners. For temperatures below  $\mathcal{O}(\text{MeV})$ ,  $g_* \sim 3.36$ .

the four dimensional Planck mass. However, in the latter case the cross-section could be actually suppressed by the supersymmetric breaking mass scale  $M_s^2$  which determines the mass of the gravitino  $m_{3/2} \sim M_s^2/M_p$ , provided  $m_g \neq m_{3/2}$ . In order to study the gravitino abundance we need to study the Boltzmann equation for the gravitino number density [21]

$$\frac{dn_{3/2}}{dt} + 3Hn_{3/2} = \langle \Sigma_{\text{tot}} v_{\text{rel}} \rangle n_{\text{rad}}^2 - \frac{m_{3/2}}{\langle E_{3/2} \rangle} \frac{n_{3/2}}{\tau_{3/2}}, \quad (14)$$

where  $\langle \dots \rangle$  represents thermal average,  $n_{\text{rad}}$  is the number density of relativistic particles  $n_{\text{rad}} \propto T^3$ ,  $v_{\text{rel}}$  is the relative velocity of the scattering radiation which in our case  $\langle v_{\text{rel}} \rangle = 1$ , and the factor  $m_{3/2}/\langle E_{3/2} \rangle$  is the average Lorentz factor. We notice that in the radiation era non-conventional brane cosmology gives the following Hubble expansion

$$H \approx \left( \frac{4\pi^5}{3} \right)^{1/2} \frac{g_*}{30} \frac{T^4}{\sqrt{\lambda} M_p}. \quad (15)$$

In supersymmetric version  $g_* \sim 300$  if the reheat temperature is more than the masses of the superpartners. It is worth mentioning that the scale factor during the radiation era follows  $a(t) \propto t^{1/4}$ , which is contrary to the standard Big Bang scenario where  $a(t) \propto t^{1/2}$ . However, we must not forget that the derivation is based on the fact that we are in a regime where  $\rho > 2\lambda$ . In Eq. (14), after the end of inflation the first term in the right-hand side dominates the second. If we assume the adiabatic expansion of the Universe  $a \propto t^{-1}$ , then we can rewrite Eq. (14) as  $Y_{3/2} = (n_{3/2}/n_{\text{rad}})$ .

$$\frac{dY_{3/2}}{dT} \approx - \frac{\langle \Sigma_{\text{tot}} \rangle n_{\text{rad}}}{HT}. \quad (16)$$

We notice that we can integrate the temperature dependence from this equation and the expression is almost independent of the temperature. we mention here that the above expression is exactly the same as in the standard Big Bang case [21]. However, Eq. (16) does not produce the correct value for  $Y_{3/2}$ , since the true conserved quantity is the entropy per comoving volume. In our case if we assume the gravitinos do not decay within the time frame we are interested in, then we may be able to get an expression for the gravitino abundance at two different temperatures

$$Y_{3/2}(T) \approx \frac{g_*(T)}{g_*(T_{\text{rh}})} \frac{n_{\text{rad}}(T_{\text{rh}}) \langle \Sigma_{\text{tot}} \rangle}{H(T_{\text{rh}})}. \quad (17)$$

Here we assume that the initial abundance of gravitinos at  $T_{\text{rh}}$  is known to us, and the dilution factor  $g_*(T)/g_*(T_{\text{rh}})$  takes care of the decrease in the relativistic degrees of freedom. For a rough estimate we assume the total cross-section  $\Sigma_{\text{tot}} \propto 1/M_p^2$ , and,  $n_{\text{rad}}(T_{\text{rh}}) \propto T_{\text{rh}}^3$ , we finally get an expression for the gravitino abundance at temperature  $T$  [18]

$$Y_{3/2}(T < \text{MeV}) \approx 10^{-3} \frac{\sqrt{\lambda}}{T_{\text{rh}} M_p}, \quad (18)$$

where we have assumed the dilution factor to be  $\sim 10^{-2}$ . If we assume that the inflaton decays via the Yukawa couplings, then with the help of Eq. (8) we get a simple expression

$$Y_{3/2}(T) \approx 10^{-7/4} \frac{\lambda^{1/3}}{M_p^{4/3}} \left( \frac{g_*}{\alpha_\phi} \right)^{1/4}, \\ \approx 10^{-7/4} \left( \frac{M_5}{M_p} \right)^2 \left( \frac{g_*}{\alpha_\phi} \right)^{1/4}, \quad (19)$$

where  $g_*$  is actually evaluated at  $T_{\text{rh}}$ , and, the abundance expression is true for temperatures below MeV. The above expression for the abundance of the gravitinos is quite important. It tells us directly that higher the value of  $M_5$  is, higher the abundance is. For an example if  $M_5 \sim 10^{15}$  GeV, the abundance is roughly  $Y_{3/2} \approx 10^{-8}$ , which is more than the required acceptable bound  $Y_{3/2} \leq 10^{-10}$  [25] for a successful nucleosynthesis. This puts severe constraint on the five dimensional Planck mass

$$M_5 \leq 10^{-5} M_p, \quad (20)$$

where we have taken  $\alpha \sim 10^{-1}$ ,  $g_* \sim 300$ , and the brane tension is given by

$$\lambda \leq 10^{-31} M_p^4. \quad (21)$$

Thus we see that introducing supersymmetry in the brane leads to lowering the upper bound on the five dimensional Planck mass and also the brane tension in the brane-world set-up.

Following Eq. (14) we notice that as the Universe expands the Hubble time  $H^{-1} \rightarrow \tau_{3/2}$ , and at the time of decay the last term starts dominating the rest. Setting  $M_{3/2}/\langle E_{3/2} \rangle = 1$ , it is then easy to estimate the gravitino abundance, which yields

$$Y_{3/2}(t) \approx 10^{-3} \frac{\sqrt{\lambda}}{T_{\text{rh}} M_p} e^{-t/\tau_{3/2}}, \quad (22)$$

where  $\tau_{3/2}$  is the life time of the gravitino. We notice that as the five dimensional Planck mass  $M_5$  increases, the transition to the conventional cosmology becomes swifter. This suggests that the gravitinos with mass  $\sim \text{TeV}$  would decay close to  $T_{\text{decay}} \approx \mathcal{O}(\text{MeV})$ . This is because the Universe follows the conventional cosmology with  $t \sim (T/\text{MeV})^{-2} \text{sec}$  after the transition. On the other hand if  $M_5$  is smaller then one could expect that the temperature at which the gravitinos decay would be higher, determined by Eq. (15). However, one could also confirm that throughout the decay life time the evolution of the Universe would not be the same as determined by Eq. (15), because of the transition from  $a \propto t^{1/4}$  to  $a \propto t^{1/2}$  taking place before the gravitinos could decay. If

this is so then it is possible to roughly estimate the temperature at which the gravitinos would decay denoted by  $T_{\text{decay}}$ .

$$\frac{\tau_{3/2}}{\tau_{\text{rh}}} \approx \frac{T_{\text{rh}}^4}{T_{\text{decay}}^2 T_{\text{transit}}^2}, \quad (23)$$

where  $\tau_{\text{rh}}$  can be assumed to be the time when the gravitinos are created in a thermal bath, which can be estimated from Eq. (15). With the help of Eqs. (8), (12) and (15), it can be shown that the temperature at which the gravitinos decay is almost insensitive to the five dimensional Planck mass

$$\begin{aligned} T_{\text{decay}} &\approx 10^{-25/2} \left( \frac{\tau_{3/2}}{\text{sec}} \right)^{-1/2} \left( \frac{M_{\text{p}}}{\text{GeV}} \right)^{1/2} \text{GeV}, \\ &\approx 10^{-3} \left( \frac{\tau_{3/2}}{\text{sec}} \right)^{-1/2} \text{GeV}. \end{aligned} \quad (24)$$

For a TeV mass gravitino, whose life time is around  $\tau_{3/2} \sim 10^4 \text{sec}$ , which leads to the decay temperature  $T_{\text{decay}} \approx 0.01 \text{MeV}$ . Thus the gravitinos decay very late and even if the five dimensional Planck mass is small they could be potentially threatening to nucleosynthesis. The only way to save nucleosynthesis is to have small gravitino abundance, which is possible only if we lower  $M_5$  considerably according to Eq. (20). The gravitinos are important from several points of view. Their decay products would always contain the lightest supersymmetric particles and they could still survive in the form of cold dark matter. Usually while the gravitinos decay they generate an entropy as discussed earlier and they could as well wash away previously obtained baryon asymmetry in the Universe. The gravitinos decaying via CP non-conserving interactions as discussed in Ref. [22] could be an interesting scenario as to regenerate baryogenesis.

### A. Non-perturbative aspects of the gravitino production

So far we have been assuming that the gravitinos have been created in a thermal bath at a temperature close to the reheat temperature. However, it has been very recently realized that the gravitinos like other particles could as well be created non-perturbatively [23]. Though the authors concentrated only upon the helicity  $\pm 3/2$  case, soon after that the mechanism for exciting the other half  $\pm 1/2$  were discussed in Ref. [24]. It was pointed out that helicity  $\pm 1/2$  gravitinos were created more abundantly than the  $\pm 3/2$  case, because helicity  $\pm 1/2$  gravitinos essentially eat the Goldstino mass, which for a single chiral field is nothing but the supersymmetric partner of the inflaton, usually called as inflatino. The inflatino mass is not suppressed by the four dimensional Planck mass and as a result the time varying mass contributing to the Goldstino could boost the production.

In the brane-world scenario we suspect that the non-perturbative production of the gravitinos could be even more important.

To keep this discussion general we discuss some of the key features of preheating. The elaborate idea has been discussed in Refs. [26]. After the end of inflation the scalar field begins oscillating around the bottom of the potential when mass of the inflaton is comparable to the Hubble expansion  $m \sim H$ . The energy density of the inflaton field  $\phi$  decreasing in a same way as a non-relativistic particle of mass  $m$ , where  $\rho_\phi = \dot{\phi}^2/2 + m^2\phi^2/2 \sim a^{-3}$ , and  $a$  is the scale factor of the Universe. During the homogeneous oscillations of the scalar field the Universe acts as a matter dominated era where the scale factor grows like  $a(t) \approx a_i(t/t_i)^{2/3}$  in the standard cosmology. In non-conventional cosmology from Eq. (6), it is clear that the scale factor would grow as  $a(t) \approx a_i(t/t_i)^{1/3}$ . Thus in a non-conventional cosmology the scale factor grows slowly compared to that in the standard case, and this has lots of interesting implications. Here we briefly sketch some of them. It is clear that the oscillations in  $\phi$  are sinusoidal with a decreasing amplitude  $\phi(t) \sim (1/m\sqrt{t})$  in the non-conventional scenario. However, in the standard cosmology the amplitude of the oscillations follows  $\phi(t) \sim (1/mt)$ , thus the oscillations die down faster in the conventional case compared to the non-conventional. For non-perturbative creation of particles it is important to have a very high amplitude oscillations. It is worth mentioning that one of the criterion for shutting off non-perturbative production is via decaying amplitude of the oscillations.

In the discussion of non-perturbative creation of particles, there is a very useful time varying parameter, known as “ $q$ ” parameter, given by  $q = \alpha_\phi^2 \phi(t)^2/m^2$ , where  $\alpha_\phi$  is the Yukawa coupling. The production of either bosons and fermions depend very crucially upon this parameter. For the bosons the occupation number for a given momentum mode  $k$  grows as  $n_k(t) \propto e^{2\mu_k mt}$ , where  $\mu_k = ((q/2)^2 - (2k/m - 1)^2)^{1/2}$  is known as the Floquet index. In non-conventional case it is possible to have  $q$  parameter bigger compared to that in the standard case. Another important factor is redshifting of the momentum, in the case of non-conventional cosmology this effect is again weaker compared to that of the conventional case due to slower expansion rate. A similar argument can also be given for fermion creation. It has been noticed that the massive fermion production rate follows  $\rho_{\text{fermion}} \propto q$  [27]. Thus clearly showing the importance of  $q$  parameter.

The general argument behind production of the gravitinos is more or less the same as any other fermions. However, it has been realized that the production of particles does depend on a specific type of inflationary model. The discussion here primarily based on inflaton being massive and derived from a superpotential  $W = \lambda\Phi^3$ . In this case it has been noticed that the number density of the gravitino can be  $n_{3/2} \sim H_i^3$ , where  $H_i$

is the Hubble parameter during inflation [24]. In non-conventional case  $H_i$  would be larger than that of the conventional case and naively one would expect enhanced production of gravitinos. However, there exists other interesting superpotentials such as  $W = \lambda\Phi(\chi^2 - \chi_0^2)$ , which leads to a potential similar to hybrid inflationary model [28], where the production of gravitinos has been noted to be  $n_{3/2} \sim (2\lambda\chi_0)^3$ , where  $2\lambda\chi_0$  represents the effective mass of the oscillating field [29]. One could as well find the abundance of the gravitinos in this case, which is given by  $n_{3/2}/s \sim (\lambda/\chi_0)T_{\text{rh}}$  [29]. Thus the non-perturbative production does depend on the model parameters as well, but whatsoever be the situation, for a given inflationary model if the production is larger compared to that of the perturbative production, one would require to invoke the bounds on  $M_5$  and  $\lambda$ . There is also an interesting proposal to create the gravitinos after reheating [30], and it would be again interesting to investigate their abundance in non-conventional cosmology.

## V. VIABILITY OF AFFLECK-DINE BARYOGENESIS IN NON-CONVENTIONAL COSMOLOGY

In this section we discuss briefly baryogenesis in the context of supersymmetry. This discussion is mainly to illustrate that the physics of the early Universe could be different in the brane-world cosmology. Our discussion will be mainly based on Refs. [36,37,18], but here we will make some additional comments. As we know that there are three main requirements for producing net baryon asymmetry in the Universe, baryon number violating interactions, C and CP violations and a departure from thermal equilibrium [31]. Grand Unified Theories (GUT) predict baryon number violation interaction at tree level and decay of a massive Higgs bosons in out of equilibrium could give rise to the observed baryon asymmetry which is roughly one part in  $10^{10}$ . Similarly in the electro-weak baryogenesis it is possible that nonperturbative effects could give rise to processes which preserve  $B - L$ , however, violates  $B + L$ , where  $B$  represents the baryon asymmetry and  $L$  represents the lepton asymmetry. It is however possible to get baryon asymmetry from the lepton asymmetry [32]. Thus there exists many ways to extract baryon asymmetry and the expansion of the Universe fulfills the third criteria of out-of-equilibrium decay. One could imagine this to happen during preheating itself and one such example has been described in Ref. [27]. In this section we would concentrate solely upon a supersymmetric mechanism for generating baryon asymmetry through the decay of sfermion condensate proposed in Ref. [33], known as Affleck-Dine (AD) mechanism. This mechanism depends crucially upon the total evolution of the AD field starting from inflation till the era when supersymmetry breaking effects become important and in our case it could be an interesting example to study the difference between the standard and

the non-conventional cosmology.

In the original AD scenario it was assumed that initially the sfermions would have to have large vacuum expectation values along the flat directions of the scalar potential. However, the flatness could be spoiled if there were bout of inflation, because inflation generically leads to a mass correction of the order of  $\mathcal{O}(H^2)$  to the sfermions, this is particularly true for  $F$ -type supersymmetric inflationary models. However, accidental cancellations could occur in the inflaton sector due to a special choice of superpotential [34], which would prevent effective mass gaining to the inflaton. The second potential threat comes from non-renormalizable terms in the superpotential, which would inevitably lift the flat directions. The two points which we have mentioned here have their bad consequences to bringing down the AD field very close to the minimum of the potential and as a result preventing the AD baryogenesis. If  $\mathcal{O}(H^2)$  correction to the mass of the AD field is negative by a choice of Kahler potential [35], then the AD field could sit at a minimum during inflation which would be different from zero. In this section we would consider a simple toy model and estimate the initial value of the AD field.

After the end of inflation eventually the bare mass term for the AD field  $\sim \mathcal{O}(m_{3/2})$  dominates the four dimensional Planck mass suppressed non-renormalizable corrections and begins oscillating when  $H \sim \tilde{m} \equiv m_{3/2}$ , where we have denoted  $\tilde{m}$  as a mass of the AD field. For our purpose we will simply assume  $\rho_\psi \approx \tilde{m}^2\psi^2$ . However, we must mention that our case is quite complicated. The reason is following. An important condition to realize the AD baryogenesis is that the thermalization due to the decay products of the inflaton field must take place after the decay of the AD field, and,  $\rho_{\text{r}\phi} > \rho_\psi$ , where  $\rho_{\text{r}\phi}$  is the energy density in radiation after the inflaton decay. This is a very stringent condition otherwise whatever baryon asymmetry generated prior to thermalization would be washed away. To prevent this happening, the decay rate of the inflaton should be sufficiently small. This tells us that if the inflaton decays much earlier via the Yukawa couplings then there is no way we can realize the AD baryogenesis. Thus, the inflaton must decay via the gravitational couplings, which is quite slow enough. However, as we have learnt in our earlier discussions that if the inflaton field decays via the gravitational coupling then after thermalization the brane tension would dominate the energy density of the thermalized plasma and the Universe would behave as if it were in the standard case without any non-conventional term in the evolution equation. This means that if we begin with energy density in the inflaton field more than the brane tension, then while the inflaton is oscillating there is a transition from non-conventional cosmology to the standard cosmology. We assume that the transition takes place instantly and this happens when  $\rho_\phi \approx m_\phi^2\phi^2(a_\phi/a)^3 \sim \lambda$  at

$$a = a_\lambda = \left( \frac{m_\phi^2 \phi^2}{\lambda} \right)^{1/3} a_\phi, \quad (25)$$

where  $a_\phi$  is the scale factor at the time when the inflaton begins oscillations,  $a$  is just the scale factor and  $m_\phi$  denotes the mass of the inflaton. We picturize a situation where the Universe began with a non-conventional cosmology, then after the end of inflation the inflaton begins oscillating, but the Universe is still non-conventional. When the Hubble parameter drops to a value  $H \sim \tilde{m}$  the oscillations in the AD field begins and at this time also the Universe is non-conventional. However, soon after oscillations in the AD field is induced, the transition from non-conventional to the standard cosmology paves its way. Since the mass of the AD field is  $\tilde{m} \sim m_{3/2} < m_\phi$  small compared to the mass of the inflaton, the oscillations in the AD field begin after the inflaton oscillations. This can be estimated by taking  $H \sim \tilde{m}$ . Since this happens when the Universe is non-conventional;  $H \approx (m_\phi^2 \phi^2 / M_p \sqrt{\lambda}) (a_\phi / a)^3 \sim \tilde{m}$ . We can estimate the scale factor when this happens

$$a = a_\psi = \left( \frac{\sqrt{\lambda}}{M_p \tilde{m}} \right)^{1/3} a_\lambda, \quad (26)$$

where we have used Eq. (25). It can be verified easily that  $a_\psi < a_\lambda$ . However, this restricts the five dimensional Planck mass  $M_5 < 10^{14}$  GeV. <sup>‡</sup> After the AD field begins oscillations we assume that the transition takes place at a given scale factor by Eq. (25). After  $a_\lambda$  the cosmology becomes the standard one and the Hubble rate is given by  $H \propto \sqrt{\rho} / M_p$ . In our set-up the inflaton decays when the Universe is already in the standard cosmology, thus we can estimate the scale factor when this happens by equating the Hubble parameter to the decay rate of the inflaton;  $H \approx (m_\phi \phi / M_p) (a_\phi / a_\lambda)^{3/2} (a_\lambda / a)^{3/2} \sim \Gamma_\phi = (m_\phi^3 / M_p^2)$ . Notice that the decay rate of the inflaton is via the gravitational coupling. This yields

$$a = a_{d\phi} = \left( \frac{\lambda M_p^2}{m_\phi^6} \right)^{1/3} a_\lambda. \quad (27)$$

It can be verified that  $a_\phi < a_\psi < a_\lambda < a_{d\phi}$ . During the oscillations of the AD field, the energy density decreases in the same fashion as in the case of inflaton. We can estimate the energy density in the AD field by

$$\begin{aligned} \rho_\psi &= \tilde{m}^2 \psi_0^2 \left( \frac{a_\psi}{a} \right)^3 \\ &= \frac{\tilde{m} \sqrt{\lambda} \psi_0^2}{M_p} \left( \frac{a_\lambda}{a} \right)^3, \end{aligned} \quad (28)$$

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<sup>‡</sup>For numerical estimations we have assumed  $\tilde{m} \approx m_{3/2} \sim 1$  TeV.

where we have assumed Eq. (26). In order to decide which field decays first we have to compare the two decay rates. The decay rate of the AD field can be taken to be  $\Gamma_\psi \sim (\tilde{m}^3 / \psi^2)$  [37]. The value of  $\psi$  can be estimated from Eq. (28). Thus the condition that  $\Gamma_\phi / \Gamma_\psi > 1$  leads to

$$\left( \frac{a}{a_\lambda} \right)^3 < \frac{m_\phi^3 M_5^3 \psi_0^2}{M_p^4 \tilde{m}^4}. \quad (29)$$

Following the values for  $m_\phi \sim 10^{-5} M_5$ ,  $M_5 \sim 10^{13}$  GeV, and  $\psi_0 \sim 10^4 M_5$ , we get  $a / a_\lambda < 10^6$ . On the other hand for the same parameters we can also estimate  $a_{d\phi} / a_\lambda \sim 10^4$ . Notice that the initial amplitude of the AD field  $\psi_0$  is taken larger than  $\phi_{\text{COBE}}$ . This tells us that the inflaton will decay first. But an important factor is that thermalization due to the decay of the inflaton field must happen after the full decay of the AD field.

Once the Universe becomes radiation dominated, the energy density of the relativistic decay products of the inflaton can be given by

$$\begin{aligned} \rho_{r\phi} &= \frac{m_\phi^6}{M_p^2} \left( \frac{a_{d\phi}}{a_\lambda} \right)^4 \left( \frac{a_\lambda}{a} \right)^4 \\ &= \left( \frac{\lambda^{4/3} M_p^{2/3}}{m_\phi^2} \right) \left( \frac{a_\lambda}{a} \right)^4, \end{aligned} \quad (30)$$

where we have used Eq. (27). Notice that while deriving the above equation we have assumed the standard cosmology and also note that during the radiation era  $\rho \propto a^{-4}$ . In the radiation dominated era the Hubble parameter in the standard cosmology is given by

$$H = \left( \frac{\lambda^{2/3}}{m_\phi M_p^{2/3}} \right) \left( \frac{a_\lambda}{a} \right)^2. \quad (31)$$

Now we must estimate when the AD field decays, following Refs. [36] and [37] we equate  $H \sim \Gamma_\psi \equiv \tilde{m}^3 / \psi^2$ . This takes place when the scale factor is given by

$$a = a_{d\psi} = \left( \frac{\lambda^{7/6} \psi_0^2}{\tilde{m}^4 m_\phi M_p^{5/3}} \right)^{1/5} a_\lambda. \quad (32)$$

It can be verified that  $\rho_{r\phi}(a_{d\psi}) > \rho_\psi(a_{d\psi})$ . Now we have to make sure that the thermalization of the inflaton field happens after the decay of the AD field. For that we need to estimate the thermalization rate of the inflaton field. Following the arguments given in Refs. [36] and [37] we get

$$\begin{aligned} \Gamma_T &\sim n_\phi \sigma \sim m_\phi \phi^2 \left( \frac{a_\phi}{a} \right)^3 \left( \frac{\alpha^2}{m_\phi^2} \right) \left( \frac{a}{a_{d\phi}} \right)^2 \\ &\sim \alpha^2 \left( \frac{\lambda^{1/3} m_\phi}{M_p^{4/3}} \right) \left( \frac{a_\lambda}{a} \right), \end{aligned} \quad (33)$$



where  $n_\phi$  is the number density of the relativistic particles,  $\sigma$  is the cross-section and  $\alpha$  is the fine structure constant. The thermalization of the Universe occurs when  $\Gamma_T \sim H$ , where  $H$  has already been estimated in Eq. (31).

$$a_T = \alpha^{-2} \left( \frac{\lambda^{1/3} M_p^{2/3}}{m_\phi^2} \right) a_\lambda. \quad (34)$$

At this point we can also check that  $a_{d\psi} < a_T$  for  $m_\phi \sim 10^{-5} M_5$ , and  $\alpha \sim 10^{-3/2}$ . The condition is satisfied for any reasonable value of  $\psi_0$  less than the four dimensional Planck mass.

At  $a_T$  we can compute the final baryon to entropy ratio given by [33]

$$n_B = \epsilon \left( \frac{\psi_0^2}{M_G^2} \right) \frac{\rho_\psi}{\tilde{m}}, \quad (35)$$

where  $\psi_0$  is the initial amplitude of the sfermion oscillations,  $M_G$  can be assumed to be an intermediate scale, could be supersymmetric grand unification scale and  $\epsilon(\psi_0^2/M_G^2)$  is the net baryon number generated by the decay of  $\psi$ . At  $a_T$  the entropy density can be calculated with the help of Eqs. (30) and (34).

$$s = (\rho_{r\phi}(a_T))^{3/4} \approx \frac{\alpha^6 m_\phi^{9/2}}{M_p^{3/2}}, \quad (36)$$

and finally the baryon to entropy ratio can be given by

$$\frac{n_B}{s} = \frac{\epsilon \psi_0^4 m_\phi^{3/2}}{M_G^2 \sqrt{\lambda} M_p^{3/2}} \equiv \frac{\epsilon \psi_0^4 m_\phi^{3/2}}{M_G^2 M_5^3 M_p^{1/2}}. \quad (37)$$

It is noticeable that the baryon to entropy ratio does not depend on  $\tilde{m}$ . However, it does depend on the brane tension and the initial amplitude of the AD field oscillations. The last step in the above equation has been expressed in terms of the five dimensional Planck mass. For an example, we may take  $M_G \sim 10^{15}$  GeV,  $m_\phi \sim 10^{-5} M_5$ , we get an estimation of the initial amplitude of oscillations in the AD field

$$\psi_0 = \left( \frac{10^{37}}{\epsilon} \right)^{1/4} \left( \frac{M_5}{\text{GeV}} \right)^{3/8} \text{ GeV}, \quad (38)$$

where we have taken the observed baryon to entropy ratio to be  $n_B/s \sim 10^{-10}$ . It is evident that the value of  $\psi_0$  is more than  $\phi_{\text{COBE}} \approx 10^2 M_5$ . However, for smaller values of  $M_5$  the amplitude could be comparable to  $\phi_{\text{COBE}}$ . In that case, situation could be different. Here we have implicitly assumed that the AD field decays after the decay of the inflaton. For smaller values of  $\psi_0$ , the situation could be reversed, in that case the AD field would decay before the inflaton decay. In such a case, the entropy produced would be simply given by the inflaton decay and we do not have to bother about actual thermalization of the relativistic particles.

Here we would like to make remark upon the baryon to entropy ratio obtained in Ref. [18]. The treatment was slightly different there and not fully correct, because the author did not consider the transition from non-conventional to the standard cosmology, which happens while the inflaton is still oscillating. This is a crucial point which also makes it different from the AD mechanism in the standard cosmology.

## VI. $M_5$ FROM M-THEORY VACUA

So far we have been setting the value of  $M_5$  at our will without much justification. However, the five dimensional Planck mass should be fixed from the observed four dimensional Planck mass, the string scale and the size of the internal dimensions. These scales are also closely related to the strength of the gauge couplings. In this regard we will only focus upon M-theory on  $S^1/Z_2$ . As we have earlier discussed in the introduction that the field theory limit of the strongly coupled string theory (or M-theory) has been shown to be 11 dimensional supergravity compactified on a manifold with boundaries expressed as  $S^1/Z_2 \times \text{CY}$  [2,38], where  $S^1/Z_2$  is a line segment of size  $\pi r$  and CY is a Calabi-Yau manifold with a volume  $V$  in 6 dimensions. In this context the five dimensional Planck mass is given by

$$M_5 = \frac{M_{11}^3}{V^{-1/3}}, \quad (39)$$

where  $M_{11}$  is the string scale in 11 dimensions. In fact there are three cases which have been discussed in several literatures. Various concerning phenomenological issues were already discussed in this regard in Ref. [39,40]. If the string scale lies above 1 TeV then one can imagine that the high energy theory is supersymmetric and supersymmetry is broken in our brane at a scale around 1 TeV. In this regard our earlier discussions would hold true. We summarise three cases here.

*Case 1*  $\rightarrow M_{11} \sim 10^{16}$  GeV. This has been discussed in Refs. [38,42]. In order to match phenomenologically preferred values for  $\alpha_{\text{GUT}}$ ,  $M_{\text{GUT}}$  and  $M_p$ , one would require the six dimensional volume of the Calabi-Yau manifold to be  $V \sim (3 \times 10^{16} \text{ GeV})^{-6}$  and the eleventh dimensional segment to be around  $\pi r \sim (4 \times 10^{15} \text{ GeV})^{-1}$ . This automatically fixes the five dimensional Planck mass to  $M_5 \sim 10^{17}$  GeV. In fact it turns out that following the limit posed in Eq. (11), the brane tension would dominate the energy density at late times, and the Universe would evolve like in the standard case. For such a large value of  $M_5$  chaotic inflation could also be problematic, because the inflaton field would eventually take a value more than the four dimensional Planck mass and thus the non-renormalizable terms could occur and spoil the inflationary potential. If the reheating happens at a low energy scale then it is possible to escape the gravitino abundance bounds as set in Eq. (19). It is obvious then

that the transition from non-conventional to the standard cosmology happens much before the Universe thermalizes. Some of the interesting cosmological aspects have been discussed in this regard [9,41,43]. Especially in Ref. [43], the authors have considered the moduli problem in the context of heterotic M-theory. Usually moduli have large vacuum expectation values and mass of the order of the gravitino mass  $m_{3/2}$ . Cosmological problems concerning light weakly interacting particles have been discussed earlier in the context of supergravity Ref. [44] and superstrings [45,46]. It has been recognized that similar to a gravitino problem, the moduli should also decay before nucleosynthesis, which provides the lower bound on their masses, and if they do not decay then they have an upper bound on their masses to avoid closing the Universe. It is worth mentioning here that the cosmological bounds on the mass of the moduli will remain be the same as in the standard Big Bang cosmology.

*Case 2*  $\rightarrow M_{11} \geq 10^7$  GeV. This is the case when there is an upper bound on  $r$  from the experiments on gravitational forces beyond 1mm size [47,3]. In this case  $10^{16}\text{GeV} \leq M_5 \leq 10^7\text{GeV}$ . In this case we can certainly expect the deviation from the standard cosmology. The transition from non-conventional to the standard cosmology can be estimated to be happening at a temperature  $10^{15}\text{GeV} \leq T_{\text{transit}} \leq 10\text{GeV}$ . The gravitino abundance after the end of inflation would be a cause of major problem for  $10^{14}\text{GeV} \leq M_5$ . Moduli and dilaton problems would be revisited in this regime, but the problems associated to them can be ameliorated if their masses are above 1 TeV. However, their dynamics will be governed by the non-conventional cosmology.

*Case 3*  $\rightarrow M_{11} \geq 1\text{TeV}$  with  $1/r \ll 1\text{TeV}$ , this case has been discussed in Refs. [39,48]. For  $M_{11} \sim 1$  TeV, and  $V \sim 1.7$  GeV [40], the value of the five dimensional Planck mass is  $M_5 \sim 10^9$  GeV. This leads to the transition temperature  $T_{\text{transit}} \sim 10^4$  GeV. Thus this would also lead to non-conventional cosmology and the gravitino abundance would cause a problem for nucleosynthesis provided  $M_5 \geq 10^{14}$  GeV. However, for smaller values of  $M_5$  the problem could be evaded.

## VII. DISCUSSIONS

In this paper we have discussed some consequences of brane cosmology. Our inference is based upon the fact that the Friedmann equation in the early and the late cosmology could deviate from the standard one due to the presence of an extra dimension compactified on an orbifold. We have given more emphasis upon post-inflationary cosmology. We have assumed that we had a period of inflation governed by the quadratic potential whose mass is constrained by the COBE normalization. We also assume that supersymmetry is required to solve the hierarchy problem. This is quite natural if the string scale is more than TeV. In this regard we

have noticed two important points. It is not always true that post-inflationary brane cosmology will deviate from the standard cosmology. If the decay rate of the inflaton is very slow, especially if they decay via four dimensional Planck mass suppressed couplings, then no matter whatsoever be the five dimensional Planck mass is, the Universe will end up with radiation energy density much less than the brane tension. This would automatically lead to the standard cosmology. In the opposite limit, if the inflaton decays fast enough via Yukawa or gauge couplings then the gravitinos produced from the thermal bath would have abundance crucially depending upon the five dimensional Planck mass. As we know that gravitinos decay very late, thus their decay products could generate enough entropy to wash out previously obtained baryon asymmetry and also harming synthesis of light elements. This is the reason why the gravitino abundance is a cause of great concern in any cosmological set-up. In this paper we have put some constraints upon the five dimensional Planck mass from the gravitino constraints coming from nucleosynthesis. We have noticed that the Universe undergoes a transition from non-conventional cosmological evolution to the standard evolution before nucleosynthesis takes place, and it happens at a temperature which again depends upon the five dimensional Planck mass. Interestingly this transition could happen very late provided the five dimensional Planck mass is small enough. If it happens after supersymmetry breaking scale in the observable world then there could be some interesting consequences, such as extremely low abundance of gravitinos, and, realization of Affleck-Dine baryogenesis at very low temperatures. There could be many other interesting scenarios which we have not discussed here such as entropy generation from the decay of moduli could be ameliorated. Finally we studied the Affleck-Dine baryogenesis in some detail. We have noticed that such a scheme could be realizable only if the inflaton decays very slowly. If this is so then while the inflaton is oscillating and before it decays the Universe undergoes a transition from non-conventional to the standard one. This makes the overall discussion bit more complicated, but nonetheless it is possible to realize such a baryogenesis.

## VIII. ACKNOWLEDGEMENTS

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